We discuss how the dynamics of the evolving hot fireball of quark–gluon matter impacts phase transition between the deconfined and confined state of matter. The rapid expansion of the fireball of deconfined matter created in heavy ion collisions facilitates formation of an over-saturated strange quark phase space. The related excess abundance of strangeness is compensating the suppression of this semi-heavy quark yield by its quark mass. In addition, the dynamical expansion of colored quanta pushes against the vacuum structure, with a resulting supercooling of the transition temperature. We address the status of the search for the phase boundary as function of reaction energy and collision centrality and show evidence for a change in reaction mechanism at sufficiently low energies. The phase diagram derived from the study of hadron production conditions shows two boundaries, one corresponding to the expected transition between confined and deconfined matter, with a downward temperature shift, and the other a high quark density hadronization which appears to involve heavy effective quarks, at relatively large temperatures.

1. INTRODUCTION

Quark–gluon plasma phase (QGP) is the equilibrium state of deconfined hadronic matter at high temperature and/or density. It is believed that this state has been present in the early Universe, 10—20µs into its evolution. Today, we recreate the conditions of the early Universe in laboratory experiments colliding atomic nuclei (heavy ions) at highest available energies. Colliding heaviest nuclei, we explore a domain of space and time much larger than normal hadron size, in which color-charged quarks and gluons are propagating constrained by external ‘frozen vacuum’, which abhors color. For the past four years, dedicated experimental program has been carried out at the RHIC collider.

It is an open question, if within the short time, \(10^{-22}–10^{-23}\) s, available in a laboratory heavy ion collision experiment, the confined color frozen nuclear phase can melt and turn into the deconfined QGP state of matter. There is no answer available today, nor as it seems, will a first principles simulation of the dynamic heavy ion environment become available in the foreseeable future. Our attention turns to the study of QGP observables.

The pattern of production of strangeness \(s\), and more generally, strange hadrons, is an important observable of QGP; for further theoretical details and historical developments see our book. Here, we will discuss how strangeness production within a fireball of rapidly expanding matter can strengthen the phase transition between QGP and HG. This is of importance since the QCD thermal state turns out to remain below a first order transition for a physical set of masses of two low mass \(m_q/T \to 0\), and one semi-heavy quark of mass \(m_s/T \simeq 1\). This is illustrated in the \(m_s–m_q\) plane in Fig. [1].
The presence of a positive quark (baryo) chemical potential $\mu_B = 3\mu_q$ enhances the quark number, and suppresses anti-quark number. Still, the effect of $\mu_B$ is to increase the pressure, and the effect is important when $T \simeq \mu_q$. The nearly massless quarks respond more strongly to the finite baryo-chemical potential, than do massive hadrons, and that is why the quark-phase response matters more. For this reason, a 1st-order phase transition is expected to arise at a finite value of the chemical potential [4,5,6]. Baryon density, as expressed by the value of the baryo-chemical potential $\mu_B$, is relatively low at RHIC, where $\mu_B \simeq 25$ MeV. At the much higher energy at LHC, the expected value of baryo-chemical potential is of magnitude $\mu_B \simeq 1$–2 MeV [7].

We recognize in Fig. 1 that if we could compensate the effect of the finite mass effect [9], in the limit $\gamma_s \rightarrow e^{m_s/T}$ the strange quark would play a similar role as a light quark and we expect a rather strong first order phase transition even at vanishing baryo chemical potential.

The question we address next is, if such an over-saturation of the strangeness phase space near to the phase transition is possible in the dynamical environment we study. We also look at any further dynamical effects associated with the explosive flow of deconfined matter, and search to understand if this assists development of the singular phase behavior. Following these studies we explore the current RHIC and future LHC environments from this perspective and determine the phase limit between hadron and quark phases. Our study of the experimental hadronization conditions suggests further that at relatively low collision energies the final state hadrons emerge from a high baryon density phase with the mass of constituents being indicating constituent quark mass phase [8].

2. QGP PHASE OVERSATURATION

The fireball of QGP created in heavy ion collisions is initially significantly more dense and hot than after its expansion towards final breakup condition. This expansion dilutes the high strangeness yield attained in the initially very dense and hot phase. Contrary to intuition, this can result in an over-saturation of the chemical abundance, even if the initial state is practically strangeness free.

In order to understand this effect qualitatively, consider the yield of strange quark pairs in Boltzmann approximation, at a temperature $T = T_e$, time $t = t_e$, within the volume $V = V_e$ we have:

$$N_s(t_e) = \frac{2V_eT_e^3}{\pi^2} x_e^2 K_2(x_e), \quad x_e = \frac{m_s(T_e)}{T_e}$$

We have shown above that the mass of strange quark is $T$-dependent as the scale of energy at which its value is determined is in the domain where a rapid mass change occurs [2]. Moreover, as we shall argue at the end of this work, there is evidence for presence of heavy quasi-partons in...
the deconfined phase, and such an effective thermal mass is strongly temperature dependent.

We choose the value of $T_c$ to be the point where the system has nearly reached chemical equilibrium abundance in QGP, with $\gamma_e^{\text{QGP}} = 1$ (superscript QGP reminds us that we are considering the deconfined phase). We assume that the continued expansion preserves entropy as is appropriate for an ideal liquid. Since the entropy is governed by essentially massless quark–gluon quanta, this implies that $VT^3 = \text{Const.}$. Model calculations show that, for typical values of $T_c$, the change in the absolute number of strange quark pairs has essentially stopped for $T < T_c$. Then for $t_1 > t_e$:

$$\frac{N_s(t = t_1)}{N_s(t = t_e)} \simeq 1 = \gamma_s(T_1) \frac{x^2 K_2(x_1)}{x^2 K_2(x_e)}.$$

Since the function $x^2 K_2(x)$ is a monotonically falling function (see Fig. 10.1 p197 in Ref. [2]), in general $\gamma_s^{\text{QGP}}(t_1) > 1$. In the non-relativistic limit $m > T$, that is $x > 1$:

$$\gamma_s^{\text{QGP}}(T_1) = \frac{x^2 K_2(x_1)}{x^2 K_2(x_e)} > 1.$$

How large can $\gamma_s(T_1)$ be? Should the QGP breakup occur from a supercooled state with $T_1 = 140$ MeV, than it is appropriate to consider $m_s(T_1)/m_s(T_0) \simeq 1.5-2$, with $T_0 \simeq 210$ MeV. To convert from temperature dependence to the energy scale dependence, we set $m \simeq 2\pi T$ and thus for $m(T = T_0)$, we need and estimate of $m_s(\mu = 1.3$ GeV). The PDG value is $m_s(\mu = 2$ GeV) = 80–130 MeV. Thus assuming $m_s(T_0) \simeq 140$ we find $\gamma_s^{\text{QGP}}(T_1) \rightarrow 1.9$. As this example shows it is possible to nearly completely compensate the effect of the strange quark mass, see Eq. (1).

Though $\gamma_s$ is a useful tool to understand how strangeness can help to facilitate phase transition, it is important to remember that it is merely a quantity, which relates actual abundance of strangeness to the equilibrium abundance,

$$n_s \simeq \gamma_s n_s^\infty, \quad n_s^\infty \equiv n_s(\gamma_s = 1).$$

A more direct way to see how 2+1 flavors turns into a 3-flavor QCD is to note that we hope and expect that the dynamics of fireball expansion could help us to reach the condition:

$$n_s(T_1) \simeq n_{u}(T_1) \simeq n_{d}(T_1)$$

even if at point of chemical equilibrium, $n_s(T_c) \simeq 0.5 n_{u}(T_c), i = u, d$.

In physical terms relative importance of strangeness increases since strangeness yield is not reduced along with light quark and gluon yields during the dense matter expansion. One can argue this in several ways. One is to look at the buildup of collective expansion of the fireball of matter which requires conversion of thermal energy into kinetic energy of collective motion. The entropy density decreases, but the expansion assures that the total entropy remains constant, or slightly increases. As the energy is transferred into the transverse expansion, strange quark pair yield, being weaker coupled, remain least influenced by this loss which mainly consumes light quarks and gluons.

We thus conclude that in the event the initial conditions present in QGP are sufficiently extreme to generate strangeness rapidly and abundantly, one can expect over-population of strangeness in the final QGP breakup condition can facilitate the occurrence of a phase transition at small and vanishing baryon density. In fact there is strong evidence that this can happen: we have performed model calculations of strangeness (over)population at RHIC prior to the first experimental results becoming available [11]. We found $\gamma_s(T_1)^{\text{QGP}} < 1.2$.

We believe that greater values of $\gamma_s^{\text{QGP}}$ can be easily arrived at; in the model considered the instantaneous establishment of transverse expansion speed cut into the lifespan of the QGP phase and thus reduced the production yield of strangeness. In addition the quark mass chosen was significantly above the currently preferred value $(m_s = 105 \pm 25\text{MeV})$ which has further cut into the rate of production of strangeness.

However, some other studies have found considerably smaller values of $\gamma_s^{\text{QGP}}$ at RHIC [12,13,14]. In part, this is due to employment of dynamical QGP equlibration models that do not allow gluon chemical equilibration. This maybe due to lack in these approach of rapid gluon equilibration pro-
cesses, which are presently not fully understood. We assumed rather rapid glue chemical equilibration with relaxation time shorter than 1.5 fm. Moreover, we note somewhat unrealistically low values used for the coupling constant $\alpha_s$ noted – we use QCD measured value: $\alpha_s(M_Z) = 0.118$ and use 2 and higher loop evolution to obtain the low energy scale values $^{15}$. We believe that if not at RHIC, than at LHC, we should expect at QGP hadronization a substantial phase space excess of strange quarks.

3. EXPLOSIVE MATTER FLOW

Gibbs characterized in detail the boundary conditions between two phases of matter. The Gibbs condition of importance here is:

$$P = P_A - P_B = 0.$$  

(7)

This condition assures force equilibrium: the boundary of phase $A$ does not move sine phase $B$ balances it exactly. This condition is thus assuring the dynamical stability of the phase boundary between phases $A$ and $B$.

The covariant characterization of the Gibbs condition requires the introduction of the energy–momentum tensor $T^{\mu\nu}$. In the laboratory rest frame its components are:

$$\hat{T}^{ij} = P \delta_{ij}, \quad \hat{T}^{00} = 0, \quad \hat{T}^{i0} = \varepsilon.$$  

(8)

where $\varepsilon$ is the energy density. The Latin indices as usual refer to space component $i = 1, 2, 3$ and the wide hat indicates the laboratory frame.

Gibbs considered a space-like surface along which pressure difference had to vanish. This surface is invariantly characterized by a normal four-vector:

$$n^\mu = (0, \vec{n}).$$  

(9)

We take as the covariant, frame of reference independent statement of the Gibbs condition Eq. (7):

$$T^{\mu\nu}n_\mu n_\nu = T_A^{\mu\nu}n_\mu n_\nu - T_B^{\mu\nu}n_\mu n_\nu = 0.$$  

(10)

We now consider matter subject to expansion flow. $\vec{v}$ is the velocity of the local matter element, and its 4-velocity is:

$$u^\mu = (\gamma, \vec{v}\gamma), \quad \gamma = \frac{1}{\sqrt{1 - v^2}}.$$  

(11)

The natural presence of two Lorentz vectors $n^\mu, u^\mu$, assumes that we cannot transform away the effect of motion, the colored state pushes against Gibbs surface where the phase boundary is located. We recognize this as the hadronization hyper-surface, where the final state hadrons are born.

The components of interest in the energy momentum tensors are the pressure components:

$$T^{0j} = P \delta_{0j} + (P + \varepsilon) \frac{v_j v_i}{1 - v^2}. $$  

(12)

The balance of forces between the deconfined QGP and confined hadron phase comprises the effect of the vacuum which confines color. This pressure, introduced in bag models of hadrons for the first time, is traditionally referred to as the bag constant, $B \approx (0.2\text{GeV})^4$ $^{10}$. This vacuum structure can be represented within $T^{ij}$ by:

$$P_{\nu} = -B, \quad P_{\nu} + \varepsilon_{\nu} = 0.$$  

(13)

The vacuum structure component is thus not entering the dynamical flow term, the last term in Eq. (12).

We obtain from Eq. (10):

$$B = P_p + (P_p + \varepsilon_p) \frac{\kappa v^2}{1 - v^2}, \quad \kappa = \frac{(\vec{v} \cdot \vec{n})^2}{v^2}.$$  

(14)

Here,

$$P_p = P_{\text{QGP}}^p - P_{\text{HG}}^p, \quad \varepsilon_p = \varepsilon_{\text{QGP}}^p - \varepsilon_{\text{HG}}^p$$  

(15)

are the particle pressure and energy density components in the pressure and energy density, from which any vacuum term has been separated. For $\vec{v} \to 0$, the conventional Gibbs condition reemerges:

$$P_{\text{HG}} = P_{\text{QGP}} - B.$$  

(16)

Eq. (14) describes the pressure of motion of color charged matter against the vacuum structure which is pushed out as color cannot exist there $^{18}$; we can speak of color wind $^{17}$. The magnitude of the flow effect on the value of temperature of the phase boundary is relatively large. This is due to a rather large, in comparison to the pressure, energy density. In presence of a phase transition, we can expect that the discontinuity in particle energy is about 6–7 times greater.
than the discontinuity in the particle pressure in Eq. (14) — we recall that only the total pressure is continuous as expressed by the covariant Gibbs condition, Eq. (10).

As a consequence, we expect a significant down-shift in the critical temperature: the moving colored fields cause the quark matter to supercool by as much as $\Delta T = 25 \text{ MeV}$ [18]. This means that if the equilibrium phase crossover were to occur at $T = 165 \text{ MeV}$, the dynamical cross over would be postponed to the value $T = 140 \text{ MeV}$. We do not know at this time if the collective and rapid outflow of QCD can turn a cross-over into a phase-transition.

4. STRANGENESS AND ENTROPY

The study of strangeness production requires a model of particle production in heavy ion collision. To describe the yields of particles produced we employ the statistical hadronization model (SHM) as implemented in the SHARE program package [19]. SHM is by definition a model of particle production in which the birth process of each particle fully saturates (maximizes) the quantum mechanical probability amplitude, and thus, the relative yields are determined solely by the appropriate integrals of the accessible phase space. For a system subject to global dynamical evolution such as collective flow, this is understood to apply within each local co-moving frame. Empirical evidence shows that the full SHM is capable to describe the experimental hadron particle data rather precisely.

When particles are produced in hadronization, we speak of chemical freeze-out. In order to arrive at an adequately precise understanding of the entropy production contained in the final state in the overall hadronic particle yield, one has to be sure to include all the hadronic resonances which decay feeding into the individual yield considered, e.g., the decay $K^* \rightarrow K + \pi$ feeds into $K$ and $\pi$ yields. It is important, in this context, to note that the magnitude of resonance production is sensitive to the (freeze-out) temperature at which these particles are formed.

The indirect resonance decay contribution, to particle yields, is dominant for the case of the pion yield at all hadronization conditions considered in literature. This happens even though each resonance contributes relatively little in the final count. However, the large number of resonances which contribute compensates and the sum of small contributions competes with the direct pion yield. For the more heavy hadrons, generally there is a dominant contribution from just a few particles.

In order to study the properties of the phase of matter that was created early on in the reaction, we study strangeness and entropy production. Entropy is an observable of similar character as strangeness, it is preserved or slightly increasing in the evolution of the reaction system. In fact, it is a ‘deeper’ observable than strangeness, which is produced by reactions occurring mostly after formation of entropy has occurred. Entropy enhancement, observed in terms of specific (per electrical charge) hadron multiplicity, has been perhaps the first indication of the new physics reach of CERN-SPS experimental heavy ion program [20].

Hadron formation from QGP phase has to absorb the high entropy content of QGP which originates in broken color bonds. The lightest hadron is pion and most entropy per energy is consumed in hadronization by producing these particles abundantly. It is thus important to free the yield of these particles from the chemical equilibrium constraint.

At RHIC, we have very rich data field with yields of many final state particles available. This allows us to consider both, the yield of strangeness per (nearly conserved) entropy and per (exactly conserved) baryon number. In addition, we look at the cost in final state thermal energy to make a strange quark pair.

In the QGP, the dominant entropy production occurs during the initial glue thermalization, and the thermal strangeness production occurs in parallel and/or just a short time later [21]. The entropy production occurs predominantly early on in the collision during the thermalization phase. Strangeness production by gluon fusion is most effective in the early, high temperature environment, however it continues to act during the evolution of the hot deconfined phase until
hadronization \[22\].

Both strangeness and entropy are nearly conserved in the evolution towards hadronization and thus the final state hadronic yield analysis value for \( s/S \) is closely related to the thermal processes in the fireball at \( \tau \simeq 1-2 \) fm/c. We believe that, for reactions in which the system approaches strangeness equilibrium in the QGP phase, one can expect a prescribed ratio of strangeness per entropy, the value is basically the ratio of the QGP degrees of freedom.

We estimate the magnitude of \( s/S \) deep in the QGP phase, considering the hot stage of the reaction \[23\]. For an equilibrated non-interacting QGP phase with perturbative properties:

\[
\frac{s}{S} = \frac{(3/\pi^2)T^3(m_s/T)^2K_2(m_s/T)}{(32\pi^2/45)T^3 + n_f[(7\pi^2/15)T^3 + \mu_q^2T]},
\]

\[
= \frac{0.027}{1 + 0.054(\ln \lambda_q)^2}. \quad (17)
\]

Here, we used for the number of flavors \( n_f = 2.5 \) and \( m_s/T = 1 \). We see that the result is a slowly changing function of \( \lambda_q \); for large \( \lambda_q \simeq 4 \), we find at modest SPS energies, the value of \( s/S \) is reduced by 10%. Considering the slow dependence on \( x = m_s/T \simeq 1 \) of \( W(x) = x^2K_2(x) \) there is minor dependence on the much less variable temperature \( T \).

The dependence on the degree of chemical equilibration which dominates is easily obtained separating the different degrees of freedom, and for simplicity, we look here at the case \( \lambda_q \to 1 \):

\[
\frac{s}{S} = f(\alpha_s)0.027 \frac{\gamma_s}{0.38 \gamma_G + 0.12 \gamma_s + 0.5 \gamma_H}, \quad (18)
\]

All \( \gamma_i \) refer here to QGP phase. \( f(\alpha_s) \) is an unknown factor that accounts for the interaction effects — up to this point, we had assumed that these are canceling, with \( f(\alpha_s) \to 1 \). In any case, seen the dependence of Eq. (18) on \( \gamma_i \), we expect to see a gradual increase in \( s/S \) as the QGP source of particles approaches chemical equilibrium with increasing collision energy and/or increasing volume.

Since the ratio \( s/S \) is established early on in the reaction, the above relations, and the associated chemical conditions we considered, probe the early hot phase of the fireball. In fact, at hadron freeze-out the QGP picture used above does not apply. Gluons are likely to freeze faster than quarks and both are subject to much more complex non-perturbative behavior. However, the value of \( s/S \) is nearly preserved from the hot QGP to the final state of the reaction.

How does this simple prediction compare to experiment? Given the statistical parameters, we can evaluate the yields of particles not yet measured and obtain the rapidity yields of entropy, net baryon number, net strangeness, and thermal energy, both for the total reaction system and also for the central rapidity condition, also as function of centrality.

As is seen in the top left panel of Fig. 2, the rise of strangeness yield with centrality is faster than the rise of baryon number yield. For the most central head-on reactions, we reach, at RHIC-200, \( s/B = 9.6 \pm 1 \). In other words we make more than three strange quark pairs for every valance quark that is stopped in the central rapidity region. In the middle panel of Fig. 2, we compare strangeness with entropy production \( s/S \). We see a smooth transition as function of participant number \( A \), (on left) from a flat peripheral behavior where \( s/S \lesssim 0.02 \) to smoothly increasing \( s/S \) reaching \( s/S \simeq 0.028 \) from the most central reactions. On right, in Fig. 2, we see that the change on \( s/S \) is much more drastic as function of reaction energy. In both cases (energy and centrality dependence), the ratio \( s/S \) rises smoothly to the asymptotic value expected based on the count of quark and gluons degrees of freedom inherent in Eq. (18).

In the bottom panel of Fig. 2 on left, we see the thermal energy cost \( E/s \) of producing a pair of strange quarks as function of the size of the participating volume (i.e., \( A \)). This quantity shows a smooth drop which can be associated with transfer of thermal energy into collective transverse expansion after strangeness is produced. Thus, it seems that the cost of strangeness production is independent of reaction centrality. The result is different when we consider \( \sqrt{s_{NN}} \) dependence of this quantity, see bottom panel on right. There is a very clear change in the energy efficiency of making strangeness at the threshold energy.
Figure 2. Strangeness per net baryon $s/B$, strangeness per entropy $s/S$, and $E_{th}/s$ the thermal energy cost to make strangeness. Left: as a function of centrality, Right as function of reaction energy, adopted from [9,23]. The results on right include the central rapidity conditions at RHIC energies (dashed, blue) lines. The actual results are the symbols, the lines guide the eye. Dotted lines on left are assuming $\gamma_q = 1$ all other results allow for variation of $\gamma_q$ in the fit. Dashed lines, on left, are for central rapidity, all other results are including full coverage of phase space. See Refs. [9,23] for further detail of model and method.

5. DECONFINEMENT THRESHOLD AND PHASE STRUCTURE

Above results suggest that at, sufficiently high energy and large number of participants, the system considered has as dynamical degrees of freedom quarks and gluons. In order to strengthen the evidence for the new phase of matter, we would like to be able to find experimentally and without a complex, and thus model dependent, analysis, some evidence for a change in reaction mechanism. We can consider variation of particle yields either with the centrality of collision, which fixes the initial transverse size of the fireball, and thus its size, or the variation with reaction energy. We are looking for a sharp change in particle yields. Since these vary rapidly due to global changes in variables such as the total energy, and/or volume, the more interesting experimental variable is the ratio of hadron yields.

We can consider a variable which traces out qualitatively the observables we discussed in last section. Since the yield of $K^+$ closely follows that of strangeness $s = \bar{s}$, and the yield of $\pi^+$ is related to the total multiplicity $h$, and thus entropy $S$, the experimental observable of interest is the ratio $K^+ / \pi^+ \propto \bar{s}/\bar{d}$ yield ratio [24]. This ratio has been studied experimentally as function of $\sqrt{s_{NN}}$ [25] and a pronounced ‘horn’ structure arises at relatively low reaction energies, see Fig. 3. Moreover there seems to be a raise in this ratio after a dip at intermediate energies.

This effect is understood, in data analysis, to be due to a rather sudden modification of chemical conditions in the dense matter fireball: the rapid rise in strangeness $\bar{s}$ production below, and a rise in the anti-quark $\bar{d}$ yield above the peak. The measured $K^+ / \pi^+$ ratios are fitted well as is shown by the continuous line in Fig. 3.
The horn arises solely in the one $K^+/\pi^+$ particle ratio and in a good approximation it traces out the final state valance quark ratio $\bar{s}/\bar{d}$. In language of quark phase space occupancies $\gamma_i$ and fugacities $\lambda_i$, we have:

$$\frac{K^+}{\pi^+} \rightarrow \frac{\bar{s}}{d} = F(T) \left( \sqrt{\frac{\lambda_3}{\lambda_q}} \right)^{-1} \frac{\gamma_s}{\gamma_q}.$$  \hspace{1cm} (19)

In chemical equilibrium models $\gamma_s/\gamma_q = 1$, the isospin factor $\lambda_3$ is insignificant. Thus the horn in the $K^+/\pi^+$ ratio must arise solely from the variation in the ratio $\lambda_s/\lambda_q$ and the change in temperature $T$. $F(T)$, which describes the ratio of phase spaces, is a smooth function of $T$. Normally, one expects that $T$ increases with collision energy, hence on this ground alone, we expect an monotonic increase in the $K^+/\pi^+$ ratio as function of reaction energy.

As the collision energy rises, the increased hadron yield per baryon requires a decreasing value of $\lambda_q = e^{\mu_B/3T}$. The two chemical fugacities $\lambda_s$ and $\lambda_q$ are coupled by the condition that the strangeness is conserved. This leads to a smooth $\lambda_s/\lambda_q$. Temperature of fit is usually decreasing smoothly with increasing chemical potential. Thus, without the chemical non-equilibrium, one expects and finds a smooth behavior of the $K^+/\pi^+$ ratio. Thus, models which take chemical equilibrium as dogma fail to describe this interesting experimental result.

The chemical freeze-out conditions we have determined presents, in the $T-\mu_B$ plane, a more complex picture than naively expected, see Fig. 4. The RHIC $dN/dy$ results are to outer left. They are followed by RHIC and SPS $N_4\pi$ results. The dip corresponds to the 30 and 40 $A$GeV SPS results. The top right is the lowest 20 $A$GeV SPS and top 11.6 $A$GeV AGS energy range. To guide the eye, we have added two lines connecting the fit results. We see that the chemical freeze-out temperature $T$ rises for the two lowest reaction energies 11.6 and 20 $A$GeV to near the Hagedorn temperature, $T = 160$ MeV, of boiling hadron matter.

The shape of the hadronization boundary, shown in Fig. 4 in the $T-\mu_B$ plane, is the result of a complex interplay between the dynamics of heavy ion reaction, and the properties of both phases of matter, the inside of the fireball, and the hadron phase we observe. The dynamical effect, capable to shift the location in temperature of the expected phase boundary is due to the expansion dynamics of the fireball see section 3 and effects of chemical non-equilibrium, see section 2 for full discussion.

Considering the sudden nature of the fireball breakup seen in several observables 1, we conjecture that the hadronizing fireball, leading to $\gamma_s > \gamma_q = 1.6$, super-cools and experiences a true 1st order phase transition also at small $\mu_B$.
of the QCD deconfined parton fluid. The quark dominance is necessary to understand, e.g., how the azimuthal asymmetry \( v_2 \) varies for different particles [26].

These quarks and anti-quarks have, in principle, at that stage a significant thermal mass. The evidence for this is derived from the dimensionless variable \( E/T S \) (thermal energy divided by entropy and temperature). The energy end entropy per particle of non-relativistic and semi-relativistic classical particle gas comprising both quarks and anti-quarks is (see section 10, [24]):

\[
\frac{E}{N} \approx m + 3/2 T + \ldots ,
\]

\[
\frac{S}{N} \approx 5/2 + m/T + \ldots ,
\]

\[
\frac{E}{TS} \approx \frac{m/T + 3/2}{m/T + 5/2}.
\]

The value \( E/TS \to 0.78 \) found at large reaction energies and in most central collisions [23,44] can be understood in terms of a quark matter made of particles, with \( m \propto aT \), \( a = 2 \) which is close to what is expected based on thermal QCD [27].

6. LHC

We expect considerably more violent transverse expansion of the fireball of matter created at LHC, as compared to RHIC. The kinetic energy of this transverse motion must be taken from the thermal energy of the expanding matter. This leads to a greater local cooling and thus to a greater reduction in the number of thermal quanta. The entropy density decreases, but the expansion assures that the total entropy remains constant, or slightly increases.

As the energy is transferred into the transverse expansion, it would appear that primarily gluons are feeding the expansion dynamics, while strange quark pair yield, being weaker coupled, remain least influenced by this dynamics. This mechanism helps to increase the \( K/\pi \) ratio as reaction energy increases, see Fig. 3. It is possible that strange \( s \), and \( \bar{s} \) quark densities can rival in magnitude the light quark components and thus facilitate the phase transition. This should be seen in the detail of the distribution of particle yield.

In this context, more interesting than the \( K/\pi \) ratio enhancement would be the enhancement anomaly in strange (anti-baryon) yields. With \( \gamma_s \gg 1 \), we find that the more strange baryons and anti-baryons are more abundant than the more ‘normal’ species. Specifically of interest would be \( (\Omega^- + \Omega^+)/(h^+ + h^-) \), \( (\Xi^- + \Xi^+)/ (h^+ + h^-) \), and \( 2\phi/(h^+ + h^-) \) which show an order of magnitude shift in relative production strength. Detailed predictions for the values of these particles require considerable extrapolation of physics conditions from the RHIC to LHC domain [7].

We have so far not discussed charm over-saturation. Given the large charm quark mass, we expect that most of charm quark yield is due to first hard interactions of primary partons. However, there is non-negligible thermal production and annihilation of charm pairs. The yield of strange and light quarks, at time of hadronization, exceeds by about a factor 50 or more that of charm at central rapidity. Thus, even though charm phase space occupancy \( \gamma_c \) at hadronization may largely exceed the chemical equilibrium value \( \gamma_c \gg 1 \) given hadronization temperature, \( m_c/T \approx 10-30 \), it takes a factor \( \gamma_c \simeq (10)^{1.5} = 700 \) to compensate particle yield suppression due to the high charm mass.

Said differently, while strange quarks can compete in abundance with light quarks considering \( m_s/T \approx 1 \), and \( \gamma_s > 1 \), charm (and heavier) flavor(s) will remain suppressed, in absolute yield, at the temperatures we can make presently in laboratory experiments. Consequently, they cannot play any significant role in the dynamics of the phase transition.

7. HIGHLIGHTS

Our objective has been to show that there is a good reason to expect that the behavior of the QGP formed in heavy ion collisions can deviate in a significant manner from expectations formed in study of equilibrium thermal QCD matter. We have described two main effects, the chemical non-equilibrium of strange quarks, and the pressure of flowing color charge acting on the vacuum, which, in our opinion, are at current level of knowledge relevant to the issues considered.
Of most theoretical relevance and interest are the implications of non-equilibrium hadronization on the possible change in the location and nature of the phase boundary.

We have further argued that strangeness, and entropy, are well developed tools allowing the detailed study of hot QGP phase. A systematic study of strange hadrons fingerprints the properties of a new state of matter at point of its breakup into final state hadrons.

We have shown that it is possible to describe the ‘horn’ in the $K^+ / \pi^+$ hadron ratio within the chemical non-equilibrium statistical hadronization model. The appearance of this structure is related to a rapid change in the properties of the hadronizing matter.

REFERENCES